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Leptogenesis without violation of $B - L$

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Abstract: We study the possibility of generating the observed baryon asymmetry via leptogenesis in the decay of heavy Standard Model singlet fermions which carry lepton number, in a framework without Majorana masses above the electroweak scale. Such scenario does not contain any source of total lepton number violation besides the Standard Model sphalerons, and the baryon asymmetry is generated by the interplay of lepton flavour effects and the sphaleron decoupling in the decay epoch.

Keywords: Neutrino Physics, Beyond Standard Model.

Contents

1. Introduction

One of the most appealing and natural mechanisms for generating the tiny neutrino masses experimentally allowed is the (type I) see-saw mechanism [1]. This simple extension of the Standard Model (SM) with two or three right-handed Majorana neutrinos provides also a very attractive origin for the observed Baryon Asymmetry of the Universe, via leptogenesis [2,3]. A lepton asymmetry is dynamically generated in the out of equilibrium decay of the heavy Majorana neutrinos, and then partially converted into a baryon asymmetry due to $(B+L)$-violating non-perturbative sphaleron interactions [4].

Unfortunately, a direct test of the see-saw mechanism is in general not possible. If the Dirac neutrino masses are similar to the other SM fermion masses, as naturally expected, the Majorana neutrino masses turn out to be of order $10^8 - 10^{16}$ GeV, so the right-handed (RH) neutrinos can not be produced in the LHC or future colliders, neither lead to other observable effects such as lepton flavour violating processes. On the other hand, although the see-saw mechanism can also work with Majorana masses as low as 100 GeV, which in principle are within the energy reach of the LHC, the smallness of the light neutrino masses generically implies in this case tiny neutrino Yukawa couplings and as consequence negligible mixing with the active neutrinos, leading also to unobservable effects.

In order to obtain a large active-sterile neutrino mixing, some cancellations must occur in the light neutrino mass matrix. In this context, the small neutrino masses are not explained by a see-saw mechanism, but rather by an "inverse see-saw" [5], in which the global lepton number symmetry $U(1)_L$ is slightly broken by a small parameter, μ. In other words, the fine-tuned cancellations required in the light neutrino masses are not accidental, but due to an approximate symmetry. The smallness of the parameter μ is protected from radiative corrections (even without supersymmetry), since in the limit $\mu \to 0$ a larger symmetry is realized [6].

Much attention has been devoted recently to this possibility, both in the context of LHC phenomenology [7] and in leptogenesis [8–10]. A consequence of the slightly broken $U(1)_L$ symmetry is the existence of two strongly degenerate RH neutrinos, which combine to form a quasi-Dirac fermion. This is interesting for leptogenesis, because it leads to an enhancement of the CP asymmetry, avoiding the strong bounds which apply to hierarchical RH neutrinos [11], and allowing for successful leptogenesis at much lower temperatures, $T \sim \mathcal{O}(1\text{ TeV})$ *.

We focus here on a different scenario: we assume that the small lepton number violating effects responsible of light neutrino masses are negligible during the leptogenesis epoch, so $B - L$ is effectively conserved. This can be the case if global lepton number is broken spontaneously at a scale well below the electroweak phase transition [14] or, even if lepton number is broken at high scales, it leads to a CP asymmetry too small to account for the observed baryon asymmetry. The main difference with previous approaches is that in this framework the RH neutrinos combine exactly into Dirac fermions, and the total CP asymmetry vanishes. As a consequence, in order to generate a baryon asymmetry we have to rely on i) flavour effects and ii) sphaleron departure from thermal equilibrium during the leptogenesis epoch.

Leptogenesis without neutrino Majorana masses, so-called "Dirac leptogenesis", has already been considered in the literature [15], but in a completely different set-up. In Dirac leptogenesis global lepton number remains exactly unbroken (except for the SM sphaleron interactions), so the light neutrinos are Dirac fermions made of the SU(2) doublet ν_L and the singlet RH neutrino ν_R. Realistic models contain not only the SM plus RH neutrinos, but also additional heavy particles to generate a non-zero lepton number for left-handed particles and and an equal and opposite lepton number for right-handed particles in their CP violating decay.

This paper is organized as follows. In Sec. 2 we describe our leptogenesis framework, the CP asymmetries produced in the heavy Dirac neutrino decay and the basic requirements to generate the baryon asymmetry. In Sec. 3 we write the network of Boltzmann equations relevant for leptogenesis without Majorana masses. In Sec. 4 we present our results, both in the resonant and non-resonant regimes, and we conclude in Sec. 5.

2. The Framework

Our starting point is that above the electroweak phase transition the relevant particle contents for leptogenesis is that of the SM plus a number of SM singlet Dirac fermions N_i. Without loss of generality we can work on a basis in which N_i are mass eigenstates. In this basis the Lagrangian above the electroweak phase transition can be written as:

$$\mathcal{L} = \mathcal{L}_{\rm SM} + i\overline{N}_i \partial\!\!\!/ N_i - M_i \overline{N}_i N_i - \lambda_{\alpha i}\, \widetilde{h}^\dagger\, \overline{P_R N_i} \ell_\alpha - \lambda^*_{\alpha i} \overline{\ell}_\alpha P_R N_i \widetilde{h}, \tag{2.1}$$

*See also [12, 13] for alternative (extended) models of leptogenesis with RH Majorana neutrinos at the TeV scale and without resonant enhancement.

where α, i are family indices ($\alpha = e, \mu, \tau$ and $i = 1, 2, 3, \dots$), ℓ_α are the leptonic $SU(2)$ doublets, $h = (h^+, h^0)^T$ is the Higgs field ($\widetilde{h} = i\tau_2 h^*$, with τ_2 Pauli's second matrix) and $P_{R,L}$ are the chirality projectors.

A quantitative illustration of the proposed scenario which can also account for the low energy neutrino phenomenology can be found in the context of the inverse see-saw mechanism. In this type of models [5], the lepton sector of the Standard Model is extended with two electroweak singlet two-component leptons per generation, i.e.,

$$\ell_i = \begin{pmatrix} \nu_{Li} \\ e_{Li} \end{pmatrix}, e_{Ri}, \nu_{Ri}, s_{Li} \; . \tag{2.2}$$

In the original formulation, the singlets s_{Li} were superstring inspired E(6) singlets, in contrast to the right-handed neutrinos ν_{Ri}, which are in the spinorial representation. More recently this mechanism has also arisen in the context of left-right symmetry [16] and SO(10) unified models [17].

At zero temperature the (9×9) mass matrix of the neutral lepton sector in the ν_L, ν_R^c, s_L basis is given by

$$\mathcal{M} = \begin{pmatrix} 0 & m_D & 0 \\ m_D^T & 0 & M \\ 0 & M^T & \mu \end{pmatrix} \; , \tag{2.3}$$

where m_D and M are arbitrary 3×3 complex matrices in flavour space with

$$m_D \equiv \lambda_{\alpha i}\, v \; , \tag{2.4}$$

and $v = 174$ GeV being the Higgs vacuum expectation value. Moreover μ is a 3×3 complex symmetric matrix.

The matrix $\mathcal{M}$ can be diagonalized by a unitary transformation, leading to nine mass eigenstates: three of them correspond to the observed light neutrinos, and the other six are heavy Majorana neutrinos.

In this "inverse see-saw" scheme, assuming $m_D, \mu \ll M$ the effective Majorana mass matrix for the light neutrinos is approximately given by

$$m_\nu = m_D {M^T}^{-1} \mu M^{-1} m_D^T \; , \tag{2.5}$$

while the three pairs of two–component heavy neutrinos combine to form three quasi-Dirac fermions with masses of order M. The admixture among singlet and doublet $SU(2)$ states (and the corresponding violation of unitarity in the light lepton sector) is of order m_D/M and can be large [18–20]. This is so because although M is a large mass scale suppressing the light neutrino masses, it can be much smaller than in the standard type-I see-saw scenario since light neutrino masses in Eq. (2.5) are further suppressed by the small ratio μ/M. Thus, in this scenario, for M as low as the electroweak scale the only bounds on the Yukawa couplings are those arising from constraints on violation of weak universality, lepton flavour violating processes and collider signatures [21].

It is important to notice that in the $\mu \to 0$ limit a conserved total lepton number can be defined. This can be easily seen if, together with the standard lepton number $L_{SM} = 1$

for the SM leptons we assign a lepton number $L_N = 1$ to the singlets ν_{Ri} and s_{Li}. With this assignment the mass matrix (2.3) with $\mu = 0$ conserves $L \equiv L_{SM} + L_N$. Then, the three light neutrinos are massless Weyl particles while the six heavy neutral leptons combine exactly into three Dirac fermions, N_i, which above the electroweak scale are given by:

$$N_i = s_{Li} + \nu_{Ri} \ . \tag{2.6}$$

The smallness of the μ term can be easily understood if the total lepton number is spontaneously broken by a vacuum expectation value $\langle\sigma\rangle$, with $\mu = f\langle\sigma\rangle$ [14]. In this case, light neutrino masses are a consequence of total lepton number being broken at an energy scale much lower than the electroweak scale $\langle\sigma\rangle \ll v$, and μ vanishes exactly at the heavy neutrino decay epoch. This scenario introduces one extra scalar singlet which couples with s_L and the SM Higgs as

$$\mathcal{L}_{int} = -\frac{1}{2} f_{ij} \overline{s^c_{L_i}} \sigma s_{L_j} + \lambda |h|^2 |\sigma|^2 \ , \tag{2.7}$$

and therefore can affect our results when considered in the framework of the inverse see-saw. In principle, there is a new Dirac neutrino decay channel, $N_i \to N_j \sigma$, which could be relevant for leptogenesis; in practice, as we will see the present mechanism only works for very degenerate heavy singlets, $M_1 \simeq M_2$, therefore this channel is phase-space suppressed and our analysis will remain valid.

In this framework, if μ is effectively zero at the leptogenesis epoch, all processes conserve B and L at the perturbative level. On the other hand, the sphalerons violate $B+L_{SM}$ but conserve $B-L_{SM}$ and, since the new heavy leptons are SM singlets, they do not change L_N. Therefore the SM sphaleron processes also conserve $B-L$. In brief $B-L$ is conserved by all the interactions of the model on scales above $\langle\sigma\rangle$.

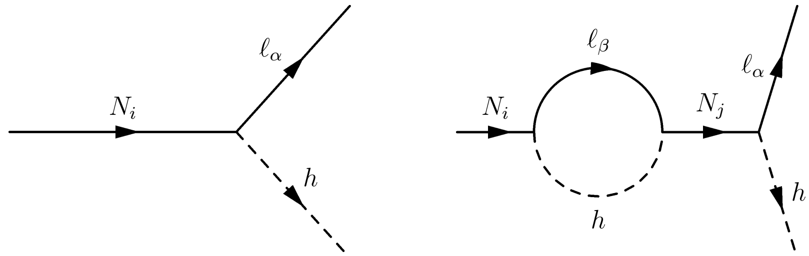
Thus, one is effectively working in the limit in which the three light neutrinos are massless, the heavy ones combine into three Dirac fermions given by Eq. (2.6) and the relevant interactions are those in Eq. (2.1).

Indeed, if leptogenesis occurs via the decay of heavy Standard Model singlets in a framework without Majorana masses above the electroweak scale, Eq. (2.1) has the relevant information. Thus our results will hold whether light neutrinos acquire masses by the inverse see-saw mechanism with the μ term generated at low scales as described above or by some other mechanism, as long as it does not imply the presence of new states relevant for leptogenesis.

2.1 The CP Asymmetries

One important ingredient that determines the baryon asymmetry generated in thermal leptogenesis in this scenario is the CP asymmetry produced in the decays of the heavy Dirac neutrinos N_i into leptons of flavour α, $\epsilon_{\alpha i}$:

$$\epsilon_{\alpha i} \equiv \frac{\Gamma(N_i \to \ell_\alpha h) - \Gamma(\bar{N}_i \to \bar{\ell}_\alpha \bar{h})}{\sum\limits_\alpha \Gamma(N_i \to \ell_\alpha h) + \Gamma(\bar{N}_i \to \bar{\ell}_\alpha \bar{h})} \ . \tag{2.8}$$



(a) Tree (b) Wave function

Figure I: The tree and one loop diagrams that contribute to the CP asymmetry in decays when the heavy neutrinos are of Dirac type.

Because of the Dirac nature of N_i, the only one-loop contribution to the CP asymmetry arises from the interference of the tree-level and the self-energy one-loop diagrams displayed in Fig. I, and it is given by [8,22] †

$$\begin{aligned}\epsilon_{\alpha i} &= \frac{-1}{8\pi(\lambda^\dagger\lambda)_{ii}} \sum_{j\neq i} \frac{a_j-1}{(a_j-1)^2+g_j^2} \mathrm{Im}\left[\lambda^*_{\alpha j}\lambda_{\alpha i}(\lambda^\dagger\lambda)_{ij}\right] \\ &= \frac{-1}{8\pi}\sum_{j\neq i} \frac{a_j-1}{(a_j-1)^2+g_j^2}(\lambda^\dagger\lambda)_{jj}\sqrt{K_{\alpha i}}\sqrt{K_{\alpha j}}\sum_{\beta\neq\alpha}\sqrt{K_{\beta i}}\sqrt{K_{\beta j}}\,p^{ij}_{\alpha\beta}\,, \end{aligned} \tag{2.9}$$

where $a_j \equiv M_j^2/M_i^2$, $g_j \equiv \Gamma_j/M_i$ and

$$\Gamma_i = \frac{M_i}{8\pi}(\lambda^\dagger\lambda)_{ii} \equiv \frac{1}{8\pi}\frac{\tilde{m}_i}{v^2}M_i^2 \tag{2.10}$$

is the decay width of N_i. In the last equation we have also introduced the effective mass $\tilde{m}_i$. Notice that the above CP asymmetry is only part of the usual wave function contribution for Majorana N_i, so that in the Dirac case the total CP asymmetry exactly vanishes

$$\epsilon_i \equiv \sum_\alpha \epsilon_{\alpha i} = 0\,, \tag{2.11}$$

by CPT invariance and unitarity. Thus in order for leptogenesis to occur we must be in a temperature regime in which flavour effects [24,25] are important.

In writing Eq. (2.9) we have expressed the Yukawa couplings in terms of the projection coefficients $K_{\alpha i}$ (related to the absolute values of the Yukawas) and some phases $\phi_{\alpha i}$ as

$$\lambda_{\alpha i} = \sqrt{K_{\alpha i}}\sqrt{(\lambda^\dagger\lambda)_{ii}}e^{i\phi_{\alpha i}}\,, \tag{2.12}$$

†We have already regularized the divergence at $M_i = M_j$ using the resummation procedure of [8]. In [23] a different calculation was performed, obtaining that the regulator of the singularity is $M_j\Gamma_j - M_i\Gamma_i$ instead of $M_i\Gamma_j$, but for the values of the widths that we consider ($\Gamma_j \gg \Gamma_i$), both results coincide.

where

$$K_{\alpha i} = \frac{\lambda_{\alpha i}\lambda^*_{\alpha i}}{(\lambda^\dagger\lambda)_{ii}} \,, \tag{2.13}$$

$$p^{ij}_{\alpha\beta} = -p^{ij}_{\beta\alpha} = -p^{ji}_{\alpha\beta} = \sin(\phi_{\alpha i} - \phi_{\alpha j} + \phi_{\beta j} - \phi_{\beta i}) \,. \tag{2.14}$$

The factor $\frac{a_j-1}{(a_j-1)^2+g_j^2}$ can be resonantly enhanced to $M_j/2\Gamma_j$ if $M_j - M_i \sim \Gamma_j/2$ [8]. In this regime the resonant contribution to the asymmetry becomes

$$\epsilon^{res}_{\alpha i} = -\frac{1}{2}\sqrt{K_{\alpha i}}\sqrt{K_{\alpha j}}\sum_{\beta\neq\alpha}\sqrt{K_{\beta i}}\sqrt{K_{\beta j}}\, p^{ij}_{\alpha\beta} \,. \tag{2.15}$$

2.2 The Generation of B: Basic Requirements

In the scenario that we are presenting $B - L_{SM} - L_N$ is conserved by all processes. Since we want to explore the possibility of generating the cosmological baryon asymmetry exclusively during the production and decay of the heavy Dirac neutrinos, we assume that at the beginning of the leptogenesis era $B - L_{SM} = 0$ and $L_N = 0$ (same abundance of N_i and $\overline{N}_i$)[‡]. Therefore it is clear that $B - L_{SM} = 0$ after the heavy neutrinos have disappeared. If the sphalerons were still active after the decay epoch, the final baryon asymmetry, being proportional to $B - L_{SM}$, would be zero.

Then, in order for leptogenesis to occur, one must be in the regime in which the sphalerons depart from equilibrium (which occurs at $T \sim T_f$) *during the decay epoch*. In this case, as described in next section, the baryon asymmetry freezes at a value $Y_B \propto Y_{B-L_{SM}}(T = T_f)$, which in general is not null.

To go on we assume that the observed baryon asymmetry is generated during the decay epoch of the lightest heavy neutrino, N_1. For the sake of concreteness we will assume that $M_3 \gg M_2 \geq M_1$ so the maximum contribution to the CP asymmetry in N_1 decays is due to N_2. Note that, in this scenario, for hierarchical masses the CP asymmetry is suppressed as $(M_1/M_2)^2$, instead of M_1/M_2 in the standard case with heavy Majorana neutrinos. This contributes to the impossibility of generating enough B with hierarchical heavy neutrinos. Moreover, if N_1 and N_2 have similar masses, they will coexist during the leptogenesis era, and there will be lepton flavour violating (but total lepton number conserving) washouts processes involving real N_2 as well as real N_1. In what follows we will refer to the washout processes involving real or virtual heavy neutrinos as lepton flavour violating washouts (LFVW).

This implies that, compared to high scale leptogenesis models with $M_i \gtrsim 10^9$ GeV, this electroweak scenario typically suffers from very strong LFVW [§]. The argument goes as follows: the intensity of the LFVW associated with processes involving real or virtual N_2

[‡]In this work we will always consider the case in which the heavy neutrinos are produced thermally exclusively through their Yukawa interactions, so that their abundance is null at the beginning of Leptogenesis. However we have also studied cases in which their initial abundance is that of a relativistic particle in equilibrium with the thermal bath (still satisfying $L_N(\text{initial}) = 0$) and found no significant differences.

[§]Similarly strong $\Delta L \neq 0$ washouts are also expected in standard (non-resonant) leptogenesis with Majorana neutrino masses at the weak scale [12,13].

is determined by the adimensional parameter $\tilde{m}_2/m_*$ [¶] . For fixed values of the Yukawas of N_2 – basically for a fixed value of the CP asymmetry in N_1 decays– $\tilde{m}_2/m_*$ scales as m_{pl}/M_2. Therefore the LFVW will be larger the lower M_2.

Consequently in order to avoid a complete erasure of the asymmetry generated in the processes involving N_1, either the CP asymmetry has to be resonantly enhanced so that the Yukawas of N_2 can be small while still having enough CP asymmetry, or some of the projectors $K_{\alpha 2}$ have to be very small to reduce the LFVW in some of the lepton flavours.

Altogether we find the following *minimum* requirements:

- The mass of N_1 must be not far from the sphaleron freeze-out temperature $M_i \lesssim \mathcal{O}(TeV)$.

- Either the CP asymmetry is resonantly enhanced or the Yukawa couplings of N_2 have a strong flavour hierarchy (i.e. $K_{\alpha 2} \ll 1$ for $\alpha = e, \mu$ or τ).

- Even in the non-resonant case there cannot be a large hierarchy between M_1 and M_2 in order to have as little suppression as possible from the $\frac{1}{a_2 - 1}$ factor.

With these requirements we conclude that to quantitatively determine the viability of this scenario we need to consider the evolution of the abundances of both N_1 and N_2 (as well as the corresponding $\overline{N}_i' s$) and the lepton flavour asymmetries. In order to do so we solve the set of Boltzmann equations which we describe next.

3. The Boltzmann Equations

In writing the relevant Boltzmann Equations we first notice that, contrary to the case in which the heavy neutrinos are of Majorana nature, the densities of N_i and $\bar{N}_i$ can be different and enter separately in the Boltzmann equations. Thus, in general there is an asymmetry between these two degrees of freedom which induces additional washout of the lepton asymmetry. Moreover, the usual $\Delta L_{SM} = 2$ processes mediated by the heavy neutrinos (like $\ell_\alpha h \to \bar{\ell}_\beta \bar{h}$, etc.) are absent, since total lepton number is perturbatively conserved. Therefore, the washout of the lepton asymmetries is due to $\Delta L_{SM} = 0$ lepton flavour violating scatterings mediated by the N_i ($\ell_\alpha h \to \ell_\beta h$, etc.), and $\Delta L_{SM} = -\Delta L_N = \pm 1$ reactions with one external N_i.

Considering all the $1 \leftrightarrow 2$ and $2 \leftrightarrow 2$ processes resulting from the Yukawa interactions of the heavy neutrinos and the Yukawa interaction of the top quark, the evolution of the

[¶]The quantity m_* is the *equilibrium mass*, which is defined by the condition $\frac{\Gamma_{N_i}}{H(T=M_i)} = \frac{\tilde{m}_i}{m_*}$, where H is the Hubble expansion rate, so that $m_* = \frac{16}{3\sqrt{5}} \pi^{5/2} \sqrt{g_{*SM}} \frac{v^2}{m_{pl}} \simeq 1,08 \times 10^{-3}$ eV (g_{*SM} is the number of Standard Model relativistic degrees of freedom at temperature T and m_{pl} is the Planck mass).

different densities is given by:

$$
\begin{aligned}
-zHs\frac{\mathrm{d}Y_{N_i+\bar{N}_i}}{\mathrm{d}z} &= \sum_\alpha \{N_i \leftrightarrow \ell_\alpha h\} + \{N_i\bar{\ell}_\alpha \leftrightarrow \bar{Q}_3 t\} + \{N_i\bar{t} \leftrightarrow \bar{Q}_3\ell_\alpha\} + \{N_i Q_3 \leftrightarrow t\ell_\alpha\} \\
&+ \sum_{\alpha,\beta,j\neq i} \{N_i\bar{N}_j \leftrightarrow \ell_\alpha\bar{\ell}_\beta\} + \{N_i\ell_\beta \leftrightarrow N_j\ell_\alpha\}' + \{N_i\bar{\ell}_\alpha \leftrightarrow N_j\bar{\ell}_\beta\} \qquad (3.1) \\
&+ \sum_{j\neq i} \{N_i\bar{N}_j \leftrightarrow h\bar{h}\} + \{N_i h \leftrightarrow N_j h\}' + \{N_i\bar{h} \leftrightarrow N_j\bar{h}\}\ , \\
-zHs\frac{\mathrm{d}Y_{N_i-\bar{N}_i}}{\mathrm{d}z} &= \sum_\alpha (N_i \leftrightarrow \ell_\alpha h) + (N_i\bar{\ell}_\alpha \leftrightarrow \bar{Q}_3 t) + (N_i\bar{t} \leftrightarrow \bar{Q}_3\ell_\alpha) + (N_i Q_3 \leftrightarrow t\ell_\alpha) \\
&+ \sum_{\alpha,\beta,j\neq i} (N_i\bar{N}_j \leftrightarrow \ell_\alpha\bar{\ell}_\beta) + (N_i\ell_\beta \leftrightarrow N_j\ell_\alpha)' + (N_i\bar{\ell}_\alpha \leftrightarrow N_j\bar{\ell}_\beta) \qquad (3.2) \\
&+ \sum_{j\neq i} (N_i\bar{N}_j \leftrightarrow h\bar{h}) + (N_i h \leftrightarrow N_j h)' + (N_i\bar{h} \leftrightarrow N_j\bar{h})\ , \\
-zHs\frac{\mathrm{d}Y_{\Delta_\alpha}}{\mathrm{d}z} &= \sum_i (N_i \leftrightarrow \ell_\alpha h) + (N_i\bar{\ell}_\alpha \leftrightarrow \bar{Q}_3 t) + (N_i\bar{t} \leftrightarrow \bar{Q}_3\ell_\alpha) + (N_i Q_3 \leftrightarrow t\ell_\alpha) \\
&+ \sum_{i,j,\beta\neq\alpha} (N_i\bar{N}_j \leftrightarrow \ell_\alpha\bar{\ell}_\beta) + (N_i\ell_\beta \leftrightarrow N_j\ell_\alpha)' + (N_i\bar{\ell}_\alpha \leftrightarrow N_j\bar{\ell}_\beta) \qquad (3.3) \\
&+ \sum_{\beta\neq\alpha} (h\bar{h} \leftrightarrow \ell_\alpha\bar{\ell}_\beta) + (\ell_\beta\bar{h} \leftrightarrow \ell_\alpha\bar{h}) + (\ell_\beta h \leftrightarrow \ell_\alpha h)'\ ,
\end{aligned}
$$

where $Y_X \equiv n_X/s$ is the number density of a single degree of freedom of the particle specie X normalized to the entropy density and $y_X \equiv (Y_X - Y_{\bar{X}})/Y_X^{eq}$ (to be used below) is the asymmetry density normalized to the equilibrium density. With Q_3 and t we denote respectively the third generation quark doublet and the top $SU(2)$ singlet. We have also defined $Y_{\Delta_\alpha} \equiv Y_B/3 - Y_{L_\alpha}$, where Y_B is the baryon asymmetry and $Y_{L_\alpha} = (2y_{\ell_\alpha} + y_{e_{R\alpha}})Y^{eq}$ is the total lepton asymmetry in the flavour α (with $Y^{eq} \equiv Y^{eq}_{\ell_\alpha} = Y^{eq}_{e_{R\alpha}}$). Moreover $Y_{N_i+\bar{N}_i} \equiv Y_{N_i} + Y_{\bar{N}_i}$ is the total normalized density of the heavy neutrino N_i and $Y_{N_i-\bar{N}_i} \equiv Y_{N_i} - Y_{\bar{N}_i}$ is the corresponding L_N asymmetry.

To write the Eqs. (3.1)–(3.3) we have also defined the following combinations of reaction densities:

$$
\begin{aligned}
[a,b,... \leftrightarrow i,j,...] &= \frac{n_a}{n_a^{eq}}\frac{n_b}{n_b^{eq}}\gamma^{eq}(a,b,... \to i,j,...) - \frac{n_i}{n_i^{eq}}\frac{n_j}{n_j^{eq}}\gamma^{eq}(i,j,... \to a,b,...), \qquad (3.4) \\
(a,b,... \leftrightarrow i,j,...) &\equiv [a,b,... \leftrightarrow i,j,...] - [\bar{a},\bar{b},... \leftrightarrow \bar{i},\bar{j},...]\ , \qquad (3.5) \\
\{a,b,... \leftrightarrow i,j,...\} &\equiv [a,b,... \leftrightarrow i,j,...] + [\bar{a},\bar{b},... \leftrightarrow \bar{i},\bar{j},...]\ , \qquad (3.6)
\end{aligned}
$$

and the prime written in the contribution of some processes indicates that the on-shell contribution to them has to be subtracted. Note that we have not included scatterings involving gauge bosons. They do not introduce qualitatively new effects and no further density asymmetries are associated to them. We have also ignored finite temperature corrections to the particle masses and couplings [26]. In particular we take all equilibrium number densities n_X^{eq}, with $X \neq N_i, \bar{N}_i$, equal to those of massless particles.

After summing over the most relevant contributions [‖] we find:

$$\frac{\mathrm{d}Y_{N_i+\bar{N}_i}}{\mathrm{d}z} = \frac{-2}{sHz}\left(\frac{Y_{N_i+\bar{N}_i}}{Y^{eq}_{N_i+\bar{N}_i}}-1\right)\sum_{\alpha}\left(\gamma^{N_i}_{\ell_\alpha h}+\gamma^{N_i\bar{\ell}_\alpha}_{\bar{Q}_3 t}+2\,\gamma^{N_i Q_3}_{t\ell_\alpha}\right)\,, \tag{3.7}$$

$$\begin{aligned}\frac{\mathrm{d}Y_{N_i-\bar{N}_i}}{\mathrm{d}z} &= \frac{-1}{sHz}\Bigg\{\sum_{\alpha}\gamma^{N_i}_{\ell_\alpha h}\left[y_{N_i}-y_{\ell_\alpha}-y_h\right]+\gamma^{N_i\bar{\ell}_\alpha}_{\bar{Q}_3 t}\left[y_{N_i}-\frac{Y_{N_i+\bar{N}_i}}{Y^{eq}_{N_i+\bar{N}_i}}y_{\ell_\alpha}+y_{Q_3}-y_t\right]\\ &\quad+\sum_{\alpha}\gamma^{N_i Q_3}_{t\ell_\alpha}\left[2y_{N_i}-2y_{\ell_\alpha}+\left(1+\frac{Y_{N_i+\bar{N}_i}}{Y^{eq}_{N_i+\bar{N}_i}}\right)(y_{Q_3}-y_t)\right]\Bigg\}\,,\end{aligned} \tag{3.8}$$

$$\begin{aligned}\frac{\mathrm{d}Y_{\Delta_\alpha}}{\mathrm{d}z} &= \frac{-1}{sHz}\Bigg\{\sum_{i}\left(\frac{Y_{N_i+\bar{N}_i}}{Y^{eq}_{N_i+\bar{N}_i}}-1\right)\epsilon_{\alpha i}\,2\sum_{\beta}\left(\gamma^{N_i}_{\ell_\beta h}+\gamma^{N_i\bar{\ell}_\beta}_{\bar{Q}_3 t}+2\,\gamma^{N_i Q_3}_{t\ell_\beta}\right)\\ &\quad+\sum_{i}\gamma^{N_i}_{\ell_\alpha h}\left[y_{N_i}-y_{\ell_\alpha}-y_h\right]+\gamma^{N_i\bar{\ell}_\alpha}_{\bar{Q}_3 t}\left[y_{N_i}-\frac{Y_{N_i+\bar{N}_i}}{Y^{eq}_{N_i+\bar{N}_i}}y_{\ell_\alpha}+y_{Q_3}-y_t\right]\\ &\quad+\sum_{i}\gamma^{N_i Q_3}_{t\ell_\alpha}\left[2y_{N_i}-2y_{\ell_\alpha}+\left(1+\frac{Y_{N_i+\bar{N}_i}}{Y^{eq}_{N_i+\bar{N}_i}}\right)(y_{Q_3}-y_t)\right]\\ &\quad+\sum_{\beta\neq\alpha}\left(\gamma^{\ell_\beta h\,\prime}_{\ell_\alpha h}+\gamma^{\ell_\beta\bar{h}}_{\ell_\alpha\bar{h}}+\gamma^{h\bar{h}}_{\ell_\alpha\bar{\ell}_\beta}\right)\left[y_{\ell_\beta}-y_{\ell_\alpha}\right]\Bigg\}\,,\end{aligned} \tag{3.9}$$

where we have introduced the notation $\gamma^{a,b,\dots}_{c,d,\dots}\equiv\gamma(a,b,\dots\to c,d,\dots)$.

In writing Eqs. (3.8) and (3.9) we have accounted for the fact that the CP asymmetry in scatterings is equal to the CP asymmetry in decays since only the wave part contributes to the CP asymmetry of the different processes (see [27]). Moreover, since the total CP asymmetries ϵ_i are null, there are no source terms proportional to $\epsilon_{\alpha i}$ in the equation for the evolution of $Y_{N_i-\bar{N}_i}$, which is therefore driven only by terms proportional to the different density asymmetries.

It is also important to notice that the equations for the asymmetries are not all independent due to the condition $\frac{\mathrm{d}Y_{B-L_{SM}-L_N}}{\mathrm{d}z}=0$, where $Y_{B-L_{SM}-L_N}\equiv Y_B-Y_{L_{SM}}-Y_{L_N}$, $Y_{L_{SM}}\equiv\sum_{\alpha}Y_{L_\alpha}$, and $Y_{L_N}\equiv\sum_i Y_{N_i-\bar{N}_i}$. If we take as initial condition that all the asymmetries are null, then $\sum_\alpha Y_{\Delta_\alpha}-\sum_i Y_{N_i-\bar{N}_i}=0$.

In Fig. II we plot the different reaction densities included in the Boltzmann equations, normalized to Hn^{eq}_{ℓ}, where H is the expansion rate of the universe. This normalization is appropriate to study the contribution of the different processes to the LFVW. We show the figure for $M_1=250$ GeV, $M_2=275$ GeV, $(\lambda^\dagger\lambda)_{11}=8.2\times10^{-15}$ ($\tilde{m}_1=10^{-3}$ eV), and $(\lambda^\dagger\lambda)_{22}=10^{-4}$, which are the values of the parameters of one of the examples given in the next section. For other values of the Yukawa couplings the reaction densities $\gamma^{N_i}_{\ell_\alpha h}$,

[‖] The scatterings mediated by the Higgs or leptons involving two external heavy neutrinos have been neglected because they are much slower than the scatterings involving only one external heavy neutrino and the top quark, since the Yukawa couplings of the heavy neutrinos are much smaller than the Yukawa of the top quark.

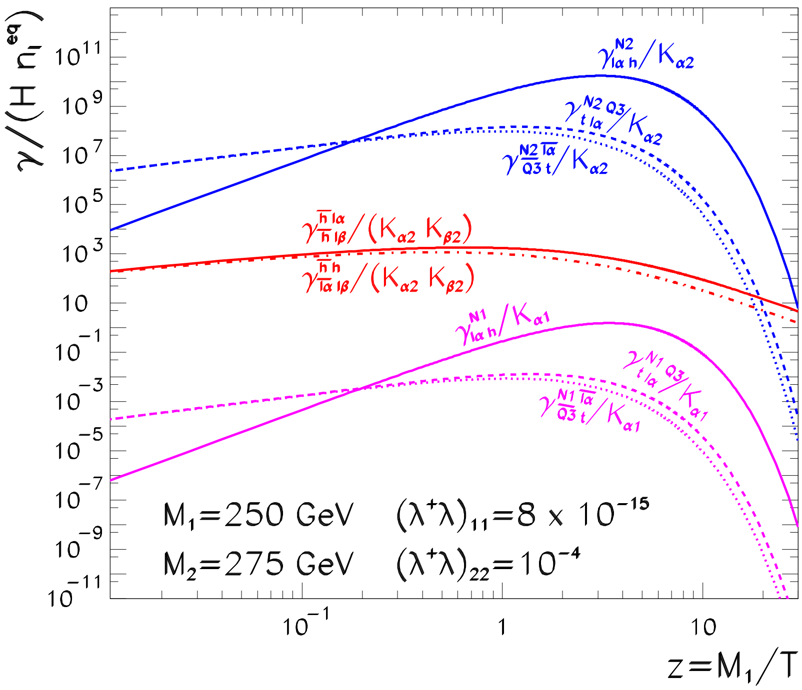


Figure II: Reaction densities included in the Boltzmann equations as a function of $z = M_1/T$, normalized to Hn_ℓ^{eq}, where H is the expansion rate of the universe.

$\gamma^{N_i \bar{\ell}_\alpha}_{\bar{Q}_3 t}$, and $\gamma^{N_i Q_3}_{t \ell_\alpha}$ scale as $(\lambda^\dagger \lambda)_{ii}$ while $\gamma^{\ell_\beta \bar{h}}_{\ell_\alpha \bar{h}}$ and $\gamma^{h\bar{h}}_{\ell_\alpha \bar{\ell}_\beta}$ scale as $((\lambda^\dagger \lambda)_{22})^2$ [**]. The subtracted s-channel reaction density $\gamma^{\ell_\beta h}_{\ell_\alpha h}{}'$ is of the same order and scales as the t- and u-channel ones shown in the plot. The infrared divergence of the reaction mediated by the Higgs in the t-channel has been regularized using a Higgs mass equal to M_1 in the propagator. For convenience we have factorized out the flavour projection factors as explicitly displayed in the figure.

From the figure we see that, as explained in Sec. 2.2, the rates of processes involving N_2 are generically very large (if the CP asymmetry in N_1 decays is required to be not too small). When $M_2 \sim M_1$ the LFVW induced by processes involving external N_2 are dominant over the ones mediated by virtual N_2 in the temperature range relevant for N_1 leptogenesis. Conversely, if there is some hierarchy between M_2 and M_1 the LFVW mediated by processes involving virtual N_2 can become the most important ones, since they are not Boltzmann suppressed for $T > M_2$.

The network of equations (3.7)–(3.9) can be solved after the densities y_{ℓ_α}, y_h and $y_t - y_{Q_3}$ are expressed in terms of the quantities Y_{Δ_α} with the help of the equilibrium conditions imposed by the fast reactions which hold in the considered temperature regime

[**] The contribution of the virtual N_1 to the LFVW processes is negligible in all the cases we are going to deal with, so we have only included N_2 virtual scatterings.

(see [28] and also [25]). For 10^6 GeV $\gg T \gtrsim T_c$ (where we denote by T_c the critical temperature of the electroweak phase transition) they read [25]

$$y_{\ell_\alpha} = -\sum_\beta C^\ell_{\alpha\beta} \frac{Y_{\Delta_\beta}}{Y^{eq}}, \qquad y_h = -\sum_\alpha C^H_\alpha \frac{Y_{\Delta_\alpha}}{Y^{eq}}, \tag{3.10}$$

with

$$C^H = \frac{8}{79}(1,\ 1,\ 1) \qquad \text{and} \qquad C^\ell = \frac{1}{711}\begin{pmatrix} 221 & -16 & -16 \\ -16 & 221 & -16 \\ -16 & -16 & 221 \end{pmatrix}. \tag{3.11}$$

Moreover, the equilibrium condition for the Yukawa interactions of the top quark implies $y_t - y_{Q_3} = y_h/2$.

Above the critical temperature, fast sphaleron processes convert the generated lepton asymmetry to baryon asymmetry [4]. Below T_c the Higgs starts to acquire its vev and this $SU(2)_L$-breaking suppresses the sphaleron rate, $\Gamma_{\Delta(B+L)}$ [29–31]. For temperatures $M_W(T) \ll T \ll M_W(T)/\alpha_W$, $\Gamma_{\Delta(B+L)} \sim M_W(M_W(T)/\alpha_W T)^3 (M_W(T)/T)^3 \exp[-E_{sp}/T]$ [8, 30] where α_W is the $SU(2)_L$ fine structure constant, $M_W(T) = g v(T)/\sqrt{2}$ is the W-boson mass and the sphaleron energy is $E_{sp} \sim M_W(T)/\alpha_W$. Because of the exponential suppression of $\Gamma_{\Delta(B+L)}$ the lepton asymmetry is not longer converted into baryon asymmetry below some temperature T_f for which $\Gamma_{\Delta(B+L)}(T_f)/H(T_f) \leq 1$.

In order to properly account for the evolution of the relevant abundances in the temperature regime $T_f < T < T_c$ one must extend the system of Boltzmann Equations to include the temperature dependent sphaleron rate. The overall effect can be approximated by replacing the usual conversion factor $n_B = \frac{28}{79} n_{(B-L_{SM})}$ by a temperature-dependent rate given by [31, 32]

$$Y_B(T) = 4\frac{77T^2 + 54v(T)^2}{869T^2 + 666v(T)^2} \sum_\alpha Y_{\Delta_\alpha}(T), \tag{3.12}$$

where $v(T)$ is the temperature-dependent Higgs vev:

$$v(T) = v\left(1 - \frac{T^2}{T_c}\right)^{\frac{1}{2}} \quad \text{with} \quad T_c = v\left(\frac{1}{4} + \frac{g'^2}{16\lambda} + \frac{3g^2}{16\lambda} + \frac{\lambda_t}{4\lambda}\right)^{-\frac{1}{2}}. \tag{3.13}$$

Here λ is the quartic Higgs self-coupling, g and g' are the $SU(2)_L$ and $U(1)_Y$ gauge couplings, and λ_t is the top Yukawa coupling.

Since the sphaleron processes are effectively switched off at $T < T_f$, the baryon asymmetry is unaffected below this temperature.

In principle an additional effect is that the set of equilibrium conditions leading to Eq. (3.11) are also modified in the temperature range between T_c and T_f. We have verified that this effect does not lead to any relevant change in our conclusions. Furthermore for $T < T_c$ the effects of the non-vanishing $v(T)$ must be accounted for in the reaction densities. As long as M_i is large enough compared to T_c these effects can be safely neglected.

4. Results

In order to quantify the required conditions for generating the observed baryon asymmetry, $Y_B = (8.75 \pm 0.23) \times 10^{-11}$ [33], we have solved the Boltzmann equations presented in the previous section. As discussed at the end of Sec. 2, the CP asymmetry in processes involving N_1 or N_2 is larger the closer the masses of N_2 and N_1 are to each other. Indeed the proximity of the masses of the heavy neutrinos is the key parameter that determines whether or not successful leptogenesis is possible. We first show how successful leptogenesis is possible in this scenario in the resonant mass regime. We then explore the requirements for obtaining the observed baryon asymmetry without reaching the resonant condition.

4.1 Resonant case

In the case that the CP asymmetry of N_1 decays receives a resonant contribution from N_2 the proposed mechanism works in a wide range of the parameter space.

As an illustration we show in Fig. III an explicit example in which

$$\begin{aligned}
M_1 &= 800 \text{ GeV} , \quad M_2 = M_1 + \frac{\Gamma_{N_2}}{2} , \\
(\lambda^\dagger\lambda)_{11} &= 10^{-12} , \quad (\lambda^\dagger\lambda)_{22} = 10^{-10} , \\
K_{e1} &= 0.3, \quad K_{\mu 1} = 0.3, \quad K_{\tau 1} = 0.4, \\
K_{e2} &= 0.1, \quad K_{\mu 2} = 0.1, \quad K_{\tau 2} = 0.8 , \\
p^{12}_{e\mu} &= p^{12}_{e\tau} = p^{12}_{\mu\tau} = 1.
\end{aligned} \tag{4.1}$$

With these values the corresponding CP asymmetries are: $\epsilon_{e1} = 6.5 \times 10^{-2}, \epsilon_{\mu 1} = 3.5 \times 10^{-2}, \epsilon_{\tau 1} = -1 \times 10^{-1}$, $\epsilon_{e2} = 1.3 \times 10^{-3}, \epsilon_{\mu 2} = 7 \times 10^{-4}$, and $\epsilon_{\tau 2} = -2 \times 10^{-3}$.

In deriving the final baryon asymmetry we have taken $m_H = 200$ GeV for which $T_c \simeq 150$ GeV and $T_f \simeq 100$ GeV.

The figure explicitly shows that $Y_{B-L_{SM}} \to 0$ as $T \to 0$ which is mandatory due to the conservation of $B - L$ in this scenario and the assumed initial conditions. Indeed it can be verified that, although Y_{Δ_α} saturate at a finite value at low temperature, with the sign assignments for the asymmetries, at any z it is verified that $Y_{B-L_{SM}} = |Y_{\Delta_e}| + |Y_{\Delta_\mu}| - |Y_{\Delta_\tau}| = Y_{L_N} = |Y_{N_1-\overline{N_1}}| + |Y_{N_2-\overline{N_2}}| \simeq |Y_{N_2-\overline{N_2}}|$. Still as illustrated in the figure, the observed baryon asymmetry is generated once the sphalerons switch off below $T_f = 100$ GeV ($z > 8$).

We notice that despite the CP asymmetry $\epsilon_{\alpha 1} \gg \epsilon_{\alpha 2}$, the N_i asymmetries verify $|Y_{N_1-\overline{N_1}}| \ll |Y_{N_2-\overline{N_2}}|$. This is so because, even though the lepton asymmetries y_{ℓ_α} are mostly produced in processes involving N_1, it is the inverse decay $\ell_\alpha h \to N_i$ what determines how much of the lepton asymmetries is transferred to the N_i asymmetry. Moreover, recall that there are no source terms proportional to $\epsilon_{\alpha i}$ in the evolution equations of $Y_{N_i-\overline{N_i}}$ Eq. (3.8). Then, since the N_2 Yukawa couplings are larger, the inverse decays of the N_2 are more efficient and a larger N_2 asymmetry is produced. Consequently for the small N_1 Yukawa couplings considered, the N_1 asymmetry plays a very little role in the dynamics of the system.

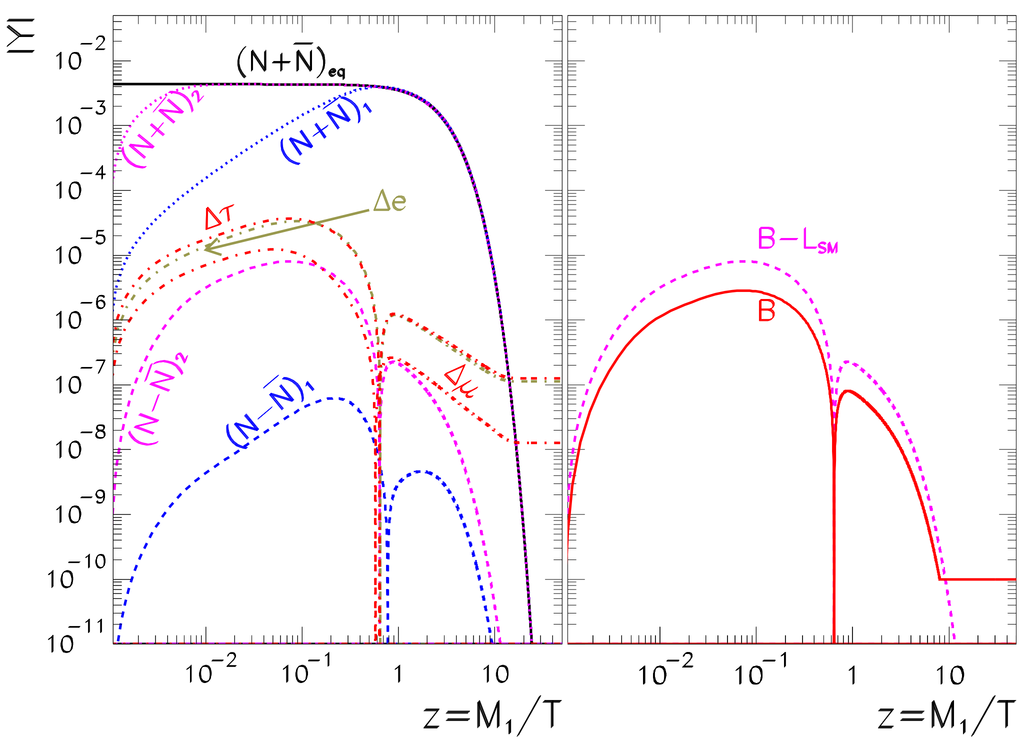


Figure III: The asymmetries $|Y_{\Delta_\alpha}|, |Y_{N_i-\bar{N}_i}|, Y_{B-L_{SM}}$, and the densities $Y_{N_i+\bar{N}_i}, Y^{eq}_{N_i+\bar{N}_i}$ as a function of z for the values of the parameters given in Eq. (4.1).

We note also that contrary to leptogenesis scenarios which occur well above the electroweak scale, once the heavy neutrinos have decayed, the universe is left with an equal amount of lepton and baryon asymmetry. Furthermore the flavour asymmetries $Y_{\Delta\alpha}$ typically remain some orders of magnitude greater than Y_B.

4.2 Non-resonant case

In this regime, in order to have enough CP asymmetry at least one of the Yukawas of N_2 has to be very large. The strongest bounds on the Yukawas of the heavy neutrinos for the range of masses we are interested in come from constraints on violation of weak universality, lepton flavour violating processes and collider signatures [21] which allow for $|(\lambda^\dagger\lambda)_{22}| \lesssim 5 \times 10^{-3}(M_2/v)^2$.

However for flavour effects to be relevant, the Yukawas of the N_2 must be smaller than the Yukawas of (at least) one of the charged leptons, so we require that $\lambda_{2\alpha} \lesssim \lambda_\tau \sim 10^{-2}$ [††]. This is so because if the fastest leptonic reaction rates were those associated to N_2, the

[††]This is a rough estimate of the validity of the "fully flavor regime" which is enough for our purposes; for a more detailed analysis about this point see [34].

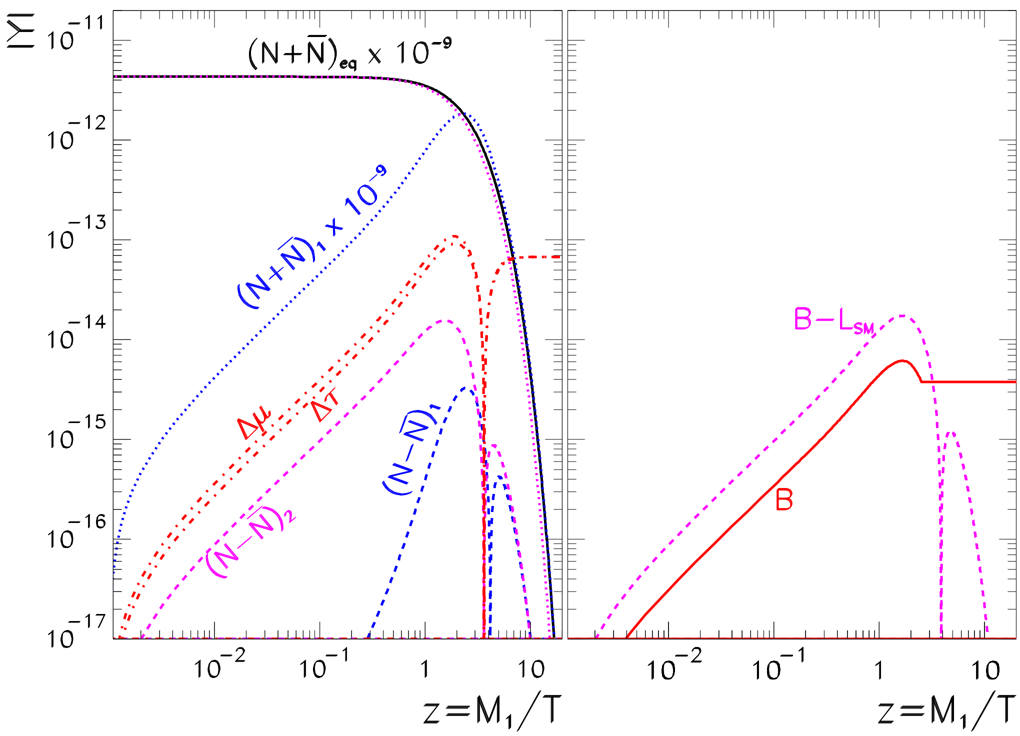


Figure IV: The asymmetries $|Y_{\Delta_\alpha}|, |Y_{N_i-\bar{N}_i}|, Y_{B-L_{SM}}$, and the densities $Y_{N_i+\bar{N}_i}, Y^{eq}_{N_i+\bar{N}_i}$ as a function of z for the values of the parameters given in Eq. (4.2). Note that the densities $Y_{N_i+\bar{N}_i}$ and $Y^{eq}_{N_i+\bar{N}_i}$ have been rescaled by 10^{-9} in order to fit into the figure.

flavour basis which diagonalizes the density matrix would be that formed by the lepton in which N_2 decays, $\ell_2 = \sum_\alpha \lambda_{\alpha 2}\ell_\alpha / \sum_\alpha |\lambda_{\alpha 2}|^2$, together with two states orthogonal to ℓ_2 and hence no baryon asymmetry would be generated since in this basis the CP asymmetry is zero.

Furthermore if the Yukawa couplings of N_2 with all the three light leptons were comparable and large (even if smaller than λ_τ) there would be strong LFVW in all flavours (see Fig. II) and the generation of the baryon asymmetry would be strongly suppressed. Therefore, as was explained before, some of the Yukawa couplings have to be very small. Note that for a fixed value of $(\lambda^\dagger\lambda)_{22}$ the CP asymmetry in a given flavour decreases linearly with decreasing Yukawa coupling of that flavour with N_2, while the washout terms decrease quadratically with it. We conclude that the most favorable situation is to have a strong hierarchy in the flavour structure of the N_2 Yukawas.

Concerning the optimum range of Yukawa couplings for N_1 there are two relevant effects. On one hand if they are too large – $\tilde{m}_1 \gtrsim m_* \simeq 10^{-3}$ eV – the maximum values of the $B - L_{SM}$ asymmetry occur at $z < 1$, so that $M_1 < T_f$ in order for the sphalerons to decouple when the baryon asymmetry is largest. In this case the analysis is more complex because one cannot neglect the effects of the breaking of $SU(2)$ in the reaction densities. The expected effect is the reduction of the N_1 decay rate due to the phase space factors. In order to determine the effect on the final B asymmetry a dedicated study is required which is beyond the scope of this paper. On the other hand, for values $\tilde{m}_1 \ll m_*$, the peak in the $B - L_{SM}$ asymmetry shifts to large values of z but it is in general lower because of the smaller production of N_1 when starting from a zero abundance as initial condition. One may wonder if this conclusion may be modified when assuming a non-vanishing initial N_1 abundance which would allow for very late N_1 decay. It is not, because at large values of z the LFVW are very suppressed and therefore the flavour effects, which are essential in this scenario, do not survive. In summary, generically larger baryon asymmetries are expected for $\tilde{m}_1 \sim m_*$.

With the above considerations in mind, we have explored the parameter space for a fixed value of M_2/M_1 (chosen near to 1) and in Fig. IV we plot the evolution of the different asymmetries and densities for a set of parameters representative of the cases with highest production of baryon asymmetry:

$$\begin{aligned} M_1 &= 250 \text{ GeV}, \quad M_2 = 275 \text{ GeV} , \\ (\lambda^\dagger\lambda)_{11} &= 8.2 \times 10^{-15} \; (\tilde{m}_1 = 10^{-3} \text{ eV}) , \\ (\lambda^\dagger\lambda)_{22} &= 10^{-4} , \\ K_{e1} &= 0. , \quad K_{\mu 1} = 0.3 , \quad K_{\tau 1} = 0.7 , \\ K_{e2} &= 0. , \quad K_{\mu 2} = 10^{-10} , \quad K_{\tau 2} \simeq 1 , \\ p^{12}_{e\mu} &= p^{12}_{e\tau} = p^{12}_{\mu\tau} = 1 . \end{aligned} \tag{4.2}$$

Since $\sqrt{(\lambda^\dagger\lambda)_{22}} > \lambda_\mu$ we have safely chosen the projectors $K_{e1}, K_{e2} = 0$ to prevent any possible flavour projection effects associated to the Yukawa interactions of N_2, which would complicate the description of the problem without substantially changing the results. In an effectively two flavour case the CP asymmetries are proportional to only one phase factor, in this case to $p^{12}_{\mu\tau}$, therefore we have adopted in the example its maximum possible value. With these values the corresponding CP asymmetries are: $\epsilon_{e1} = \epsilon_{e2} = 0, \epsilon_{\mu 1} = -\epsilon_{\tau 1} = 8.7 \times 10^{-11}, \epsilon_{\mu 2} = -\epsilon_{\tau 2} = 8.7 \times 10^{-21}$.

From the figure it can be seen that even with these large N_2 Yukawa couplings, their strong flavour hierarchy and the small $1/(a_2 - 1)$ suppression, the produced Y_B falls short to explain the observations by about 5 orders of magnitude. However, we notice that in this regime the asymmetry comes mainly from processes involving N_1. Therefore the $B - L_{SM}$ asymmetry is approximately proportional to $1/(a_2 - 1)$ which is around 5 in the example. If we take M_2 closer to M_1, that factor and the corresponding Y_B grow accordingly. Thus we see that in this regime it is also possible to generate the required baryon asymmetry as long as N_1 and N_2 are strongly degenerated even if still not in the resonant regime. For example for the values of parameters given in Eq. (4.2) the observed baryon asymmetry could be produced if $a_2 - 1 \sim 2.4 \times 10^{-5}$, which is still far from the resonance ($g_2 = 4 \times 10^{-6}$).

Something to note is that despite the great hierarchy among the projectors of N_2 onto ℓ_μ and ℓ_τ, the flavour asymmetries $Y_{\Delta\alpha}$ ($\alpha = \mu, \tau$) are quite similar in size. This is because the evolutions of the different asymmetries are strongly coupled due to the conservation of $B-L$ and the null value of the total CP asymmetry in N_1 decays.

To obtain an estimate of the order of magnitude of the density asymmetries as well as to understand their dependence on the N_2 flavour projectors (which are more relevant than the N_1 projectors when $(\lambda^\dagger\lambda)_{22} \gg (\lambda^\dagger\lambda)_{11}$), we have developed the following semiquantitative approximation:

(i) From Fig. II we see that for $K_{\mu 2} \gtrsim 10^{-9}$ the rates of the processes $N_2 \leftrightarrow \ell_{\mu,\tau} h$ are larger than the expansion rate of the Universe in the most relevant range of temperatures. Therefore at each instant the thermal bath has time to relax, i.e., the production of asymmetry equals its erasure. This implies that the derivatives of the density asymmetries are negligible with respect to the source and LFVW terms, hence we set them to zero in the Boltzmann equations.

(ii) We keep the CP asymmetries produced by N_1 (since $\epsilon_{\alpha 2} \ll \epsilon_{\alpha 1}$), as well as the N_2 decay and inverse decay reactions, and neglect all the remaining subdominant contributions, including the partial conversion of lepton asymmetry into baryon asymmetry during the leptogenesis era.

Within this approximation, Eqs. (3.8) and (3.9) simplify to

$$S_\mu(z) + K_{\mu 2}\gamma_{\ell h}^{N_2}\left[y_{N_2} - y_{\ell_\mu}\right] = 0 \,, \tag{4.3}$$

$$S_\tau(z) + K_{\tau 2}\gamma_{\ell h}^{N_2}\left[y_{N_2} - y_{\ell_\tau}\right] = 0 \,, \tag{4.4}$$

where $S_\mu(z) = -S_\tau(z) = \epsilon_{\mu 1}\left(\frac{Y_{N_1+\bar{N}_1}}{Y^{eq}_{N_1+\bar{N}_1}} - 1\right) 2\,\gamma_{\ell h}^{N_1}$ is the source term normalized to $(-sHz)^{-1}$ and $\gamma_{\ell h}^{N_i} \equiv \sum_\beta \gamma_{\ell_\beta h}^{N_i}$.

A third equation is provided by total lepton number conservation, i.e. $Y_{N_2 - \bar{N}_2} + Y_{L_\mu} + Y_{L_\tau} = 0$. In the most relevant temperature range for the model, $T \sim M_2$, we can approximate the equilibrium density of N_2 by that of one relativistic degree of freedom, therefore

$$y_{N_2} + 2y_{\ell_\mu} + 2y_{\ell_\tau} = 0 \,. \tag{4.5}$$

When $K_{\mu 2} \ll 1 \simeq K_{\tau 2}$ the solution to this system of equations is

$$\begin{aligned} y_{\ell_\mu} &= \frac{3}{5}\frac{S_\mu(z)}{K_{\mu 2}\gamma_{\ell h}^{N_2}} \,, \\ y_{N_2} = y_{\ell_\tau} &= -\frac{2}{5}\frac{S_\mu(z)}{K_{\mu 2}\gamma_{\ell h}^{N_2}} \,. \end{aligned} \tag{4.6}$$

From this analysis, it is clear that despite the large hierarchy in the projectors of N_2, all the density asymmetries have the same order of magnitude. Moreover, since the source term is proportional to $\sqrt{K_{\mu 2}}$, the density asymmetries are inversely proportional to $\sqrt{K_{\mu 2}}$. We have verified numerically that this dependency of the asymmetries on the projector actually holds in the range $10^{-9} \lesssim K_{\mu 2} \lesssim 10^{-1}$, where the approximations we made are expected to be valid.

For $\sqrt{K_{\mu 2}} \lesssim 10^{-10}$ the rates of the processes $N_2 \leftrightarrow \ell_\mu h$ are lower than the expansion rate of the Universe, hence point (i) is not longer true. In this range of $K_{\mu 2}$, the density asymmetries decrease as $\sqrt{K_{\mu 2}}$ because the main dependence on $K_{\mu 2}$ comes from the relation $\epsilon_{\mu,\tau 1} \propto \sqrt{K_{\mu 2}}$. Thus, fixing all the parameters but $K_{\mu 2}$ to the values given in Eq. (4.2), the baryon asymmetry is maximized for $K_{\mu 2}$ between 10^{-9} and 10^{-10}.

5. Summary

In this article we have studied the possibility of generating the observed baryon asymmetry via leptogenesis in the decay (and scatterings) of heavy Dirac Standard Model singlets with $\mathcal{O}$ (TeV) masses in a framework with $B - L$ conservation above the electroweak scale. In this scenario a total lepton number, which is perturbatively conserved, can be defined. This lepton number is shared between the Standard Model leptons and the heavy Dirac singlets as described in Sec.2. In this scenario, despite the total CP asymmetry is null, a CP asymmetry in the different SM lepton flavours can be generated (see Sec.2.1).

The additional physical condition for generating a non-vanishing B in this framework is that the sphalerons depart from equilibrium during the decay epoch. For symmetric initial conditions (no net baryon nor total lepton number present), $B - L_{SM} = 0$ after the heavy neutrinos have disappeared. Consequently if the sphalerons were still active after the heavy Dirac singlets decay epoch, the final baryon asymmetry, being proportional to $B - L_{SM}$, would be zero. However if they depart from equilibrium during the decay epoch the baryon asymmetry freezes at a value which in general is not null.

In summary in this scenario the baryon asymmetry is generated by the interplay of lepton flavour effects and the sphaleron decoupling in the decay epoch. In order to quantify whether enough baryon asymmetry can be generated we have constructed and solved the network of relevant Boltzmann Equations associated with the abundances of the two lightest Dirac heavy singlets, their asymmetries and the three SM flavour asymmetries. The results are given in Sec.4.

We find that the ratio of the masses of the two heavy Dirac neutrinos is the key parameter that determines whether or not successful leptogenesis is possible. The relevant Yukawa couplings are constrained from above by the requirement of having flavour effects, and from below by the requirement of large enough CP asymmetry. Within these boundaries we conclude that successful leptogenesis can occur if the masses of two heavy Dirac singlets are very degenerate $M_2/M_1 - 1 \lesssim \mathcal{O}(10^{-5})$. Recall that in our framework this degeneration is not a consequence of the small breaking of total lepton number, as in other low scale leptogenesis scenarios [8, 10], although it may be due to an additional symmetry of the heavy sector. In particular, if the CP asymmetry is resonantly enhanced – $(M_2^2 - M_1^2) \sim M_1\Gamma_2$ – , the mechanism works for a wide range of values of the Yukawa couplings and flavour projections. It is worth to explore whether the heavy neutrinos will be observable at the LHC and/or lead to measurable lepton flavour violating signals, within the parameter regions allowed by leptogenesis.

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References

[1] P. Minkowski, "Mu $\rightarrow$ E Gamma At A Rate Of One Out Of 1-Billion Muon Decays?," Phys. Lett. B **67** (1977) 421; M. Gell-Mann, P. Ramond and R. Slansky, Proceedings of the Supergravity Stony Brook Workshop, New York, 1979, eds. P. Van Nieuwenhuizen and D. Freedman (Ed. North-Holland, Amsterdam); T. Yanagida, "Horizontal symmetry and masses of neutrinos," Proceedings of the Workshop on Unified Theories and Baryon Number in the Universe, Tsukuba, Japan, 1979 (eds. A. Sawada and A. Sugamoto, KEK Report No. 79-18, Tsukuba); R. Mohapatra and G. Senjanović, "Neutrino mass and spontaneous parity nonconservation," Phys. Rev. Lett. **44** (1980) 912.

[2] M. Fukugita and T. Yanagida, "Baryogenesis Without Grand Unification," Phys. Lett. B **174** (1986) 45.

[3] S. Davidson, E. Nardi and Y. Nir, "Leptogenesis," Phys. Rept. **466** (2008) 105 [arXiv:0802.2962 [hep-ph]].

[4] V. A. Kuzmin, V. A. Rubakov and M. E. Shaposhnikov, " On the Anomalous Electroweak Baryon Number Nonconservation in the Early Universe," Phys. Lett. **B155** (1985) 36.

[5] R. N. Mohapatra and J. W. F. Valle, "Neutrino Mass and Baryon Number Nonconservation in Superstring Models," Phys. Rev. **D34** (1986) 1642.

[6] G. 't Hooft, "Naturalness, chiral symmetry, and spontaneous chiral symmetry breaking," Lecture given at Cargese Summer Inst., Cargese, France, Aug. 26 - Sep. 8, 1979.

[7] T. Han and B. Zhang, "Signatures for Majorana neutrinos at hadron colliders," Phys. Rev. Lett. **97** (2006) 171804 [arXiv:hep-ph/0604064]; F. del Aguila, J. A. Aguilar-Saavedra and R. Pittau, "Heavy neutrino signals at large hadron colliders," JHEP **0710** (2007) 047 [arXiv:hep-ph/0703261]; J. Kersten and A. Y. Smirnov, "Right-Handed Neutrinos at LHC and the Mechanism of Neutrino Mass Generation," Phys. Rev. D **76** (2007) 073005 [arXiv:0705.3221 [hep-ph]].

[8] A. Pilaftsis and T. E. J. Underwood, "Electroweak-scale resonant leptogenesis," Phys. Rev. D **72** (2005) 113001 [arXiv:hep-ph/0506107]; A. Pilaftsis, "Resonant tau leptogenesis with observable lepton number violation," Phys. Rev. Lett. **95** (2005) 081602 [arXiv:hep-ph/0408103].

[9] A. Pilaftsis, "Electroweak Resonant Leptogenesis in the Singlet Majoron Model," Phys. Rev. D **78** (2008) 013008 [arXiv:0805.1677 [hep-ph]].

[10] T. Asaka and S. Blanchet, "Leptogenesis with an almost conserved lepton number," Phys. Rev. D **78** (2008) 123527 [arXiv:0810.3015 [hep-ph]].

[11] S. Davidson and A. Ibarra, "A lower bound on the right-handed neutrino mass from leptogenesis," Phys. Lett. B **535** (2002) 25 [arXiv:hep-ph/0202239].

[12] A. Abada and M. Losada, "Leptogenesis with four gauge singlets," Nucl. Phys. B **673** (2003) 319 [arXiv:hep-ph/0306180]; A. Abada, H. Aissaoui and M. Losada, "A model for leptogenesis at the TeV scale," Nucl. Phys. B **728** (2005) 55 [arXiv:hep-ph/0409343].

[13] T. Hambye, F. S. Ling, L. Lopez Honorez and J. Rocher, "Scalar Multiplet Dark Matter," JHEP **0907** (2009) 090 [arXiv:0903.4010 [hep-ph]].

[14] M. C. Gonzalez-Garcia and J. W. F. Valle, "Fast decaying neutrinos and observable flavour violation in an new class of majoron models," Phys. Lett. B **216** (1989) 360.

[15] K. Dick, M. Lindner, M. Ratz and D. Wright, “Leptogenesis with Dirac neutrinos,” Phys. Rev. Lett. **84** (2000) 4039 [arXiv:hep-ph/9907562]; H. Murayama and A. Pierce, “Realistic Dirac leptogenesis,” Phys. Rev. Lett. **89** (2002) 271601 [arXiv:hep-ph/0206177]; B. Thomas and M. Toharia, “Phenomenology of Dirac neutrinogenesis in split supersymmetry,” Phys. Rev. D **73** (2006) 063512 [arXiv:hep-ph/0511206]; A. Bechinger and G. Seidl, “Resonant Dirac leptogenesis on throats,” [arXiv:0907.4341 [hep-ph]].

[16] E. Akhmedov, M. Lindner, E. Schnapka and J. W. F. Valle, “Dynamical Left-Right Symmetry Breaking,” Phys. Rev. D **53** (1996) 2752 [arXiv:hep-ph/9509255]; “Left-Right Symmetry Breaking In Njl Approach,” Phys. Lett. B **368** (1996) 270 [arXiv:hep-ph/9507275].

[17] S. M. Barr, “A different see-saw formula for neutrino masses,” Phys. Rev. Lett. **92** (2004) 101601 [arXiv:hep-ph/0309152]; S. M. Barr and I. Dorsner, “A prediction from the type III see-saw mechanism,” Phys. Lett. B **632** (2006) 527 [arXiv:hep-ph/0507067].

[18] J. Bernabeu, A. Santamaria, J. Vidal, A. Mendez and J. W. F. Valle, “Lepton Flavor Nonconservation at High-Energies in a Superstring Inspired Standard Model,” Phys. Lett. B **187** (1987) 303.

[19] M. C. Gonzalez-Garcia and J. W. F. Valle, “Enhanced lepton flavor violation with massless neutrinos: A Study of muon and tau decays,” Mod. Phys. Lett. A **7** (1992) 477.

[20] M. B. Gavela, T. Hambye, D. Hernandez and P. Hernandez, “Minimal Flavour Seesaw Models,” [arXiv:0906.1461 [hep-ph]].

[21] S. Antusch, C. Biggio, E. Fernandez-Martinez, M. B. Gavela and J. Lopez-Pavon, “Unitarity of the Leptonic Mixing Matrix,” JHEP **0610**, 084 (2006) [arXiv:hep-ph/0607020]; S. Antusch, J. P. Baumann and E. Fernandez-Martinez, “Non-Standard Neutrino Interactions with Matter from Physics Beyond the Standard Model,” Nucl. Phys. B **810** (2009) 369 [arXiv:0807.1003 [hep-ph]].

[22] L. Covi, E. Roulet and F. Vissani, “CP violating decays in leptogenesis scenarios,” Phys. Lett. B **384** (1996) 169 [arXiv:hep-ph/9605319].

[23] A. Anisimov, A. Broncano and M. Plumacher, “The CP-asymmetry in resonant leptogenesis,” Nucl. Phys. B **737** (2006) 176 [arXiv:hep-ph/0511248].

[24] A. Abada, S. Davidson, F. X. Josse-Michaux, M. Losada and A. Riotto, “Flavour issues in leptogenesis,” JCAP **0604** (2006) 004 [arXiv:hep-ph/0601083]; A. Abada, S. Davidson, A. Ibarra, F. X. Josse-Michaux, M. Losada and A. Riotto, “Flavour matters in leptogenesis,” arXiv:hep-ph/0605281; S. Blanchet and P. Di Bari, “Flavor effects on leptogenesis predictions,” arXiv:hep-ph/0607330.

[25] E. Nardi, Y. Nir, E. Roulet and J. Racker, “The importance of flavor in leptogenesis,” JHEP **0601** (2006) 164 [arXiv:hep-ph/0601084].

[26] G. F. Giudice *et al.*, “Towards a complete theory of thermal leptogenesis in the SM and MSSM,” Nucl. Phys. B **685** (2004) 89 [arXiv:hep-ph/0310123].

[27] E. Nardi, J. Racker and E. Roulet, “CP violation in scatterings, three body processes and the Boltzmann equations for leptogenesis,” JHEP **0709** (2007) 090 [arXiv:0707.0378 [hep-ph]].

[28] W. Buchmüller and M. Plümacher, “Spectator processes and baryogenesis” Phys. Lett. **B511** (2001) 74 [arXiv:hep-ph/0104189]; E. Nardi, Y. Nir, J. Racker, and E. Roulet, “On Higgs and sphaleron effects during the leptogenesis era”, JHEP **01** (2006) 068 [arXiv:hep-ph/0512052]

[29] L. Carson, X. Li, L. D. McLerran and R. T. Wang, " Exact computation of the small fluctuation determinant around a sphaleron" Phys. Rev. D **42** (1990) 2127.

[30] P. Arnold and L. McLerran, "Sphalerons, small fluctuations, and baryon number violation in electroweak theory," Phys. Rev. **D36** (1987) 581

[31] M. Laine and M. Shaposhnikov, "A remark on sphaleron erasure of baryon asymmetry," Phys. Rev. **D61** (2000) 117302 [arXiv:hep-ph/9911473].

[32] J. A. Harvey and M. S. Turner, "Cosmological baryon and lepton number in the presence of electroweak fermion-number violation" Phys. Rev. **D42** (1990) 3344.

[33] J. Dunkley *et al.* [WMAP Collaboration], "Five-Year Wilkinson Microwave Anisotropy Probe (WMAP) Observations: Likelihoods and Parameters from the WMAP data," Astrophys. J. Suppl. **180** (2009) 306 [arXiv:0803.0586 [astro-ph]].

[34] S. Blanchet, P. Di Bari and G. G. Raffelt, "Quantum Zeno effect and the impact of flavor in leptogenesis," JCAP **0703** (2007) 012 [arXiv:hep-ph/0611337].